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TITLE: MONTE CARLO SIMULATION OF INCLUSIVE PIONIC REACTION AROUND RESONANCE AND HIGHER ENERGY PERSPECTIVES

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Monte Carlo simulation of inclusive pionic reaction around resonance and higher energy perspectives[1]

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Abstract

We describe all the inclusive pionic reactions in a variety of nuclei around the resonance region by means of a microscopic many-body calculation that evaluates reaction probabilities, followed by a Monte Carlo simulation which follows the evolution of the pions. We also make an intrusion into the higher energy domain by paying attention to the absorption mechanisms and their repercussions in exclusive single-charge exchange reactions.

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1 Introduction

Much attention has been paid in the past to the exclusive pionic reactions in nuclei, with hopes that one can learn both about nuclear structure and reaction mechanisms. Further experimental and theoretical work has convinced us that the reaction mechanisms are more complex than were originally thought. Evidence that pion absorption requires two, three or more nucleons[1] has added new elements to the discussion. Detailed calculations also show that the amounts of two- and three-body absorption are comparable around resonance, while the four body contribution is smaller than the other two[2]. The amount of experimental data in pion nucleon reactions also calls for an unified picture that describes realistically the reaction mechanisms and which can be used to study all these reactions. We have undertaken such a task and have applied it to the study of inclusive pionic reactions where the details of nuclear structure are less important and only their gross features matter.

The study of pion propagation through a nucleus proceeds in two steps. In the first one, a theoretical scheme is described in which, by means of phenomenological couplings of mesons to nucleons and isobars, and by applying systematically many-body field theoretical methods, one is able to write down the intrinsic probabilities for a certain reaction to take place as a function of energy and the nuclear density. In the second step we translate these ideas to finite nuclei by means of a local density approximation, however taking into account finite range effects. Next, the problem of multiple collisions in the process and the interference of the different reaction channels are taken into account by means of a computer simulation of the process based upon the intrinsic reaction probabilities (for one collision) which have been calculated before. By means of such a procedure one obtains a good reproduction of the experimental cross reactions for the different reactions, absorption, quasielastic, single-charge exchange and double charge exchange, as well as differential cross sections.

We translate some of these ideas to higher energies, 400 MeV $\leq T_\pi \leq$900 MeV and indicate the importance of two- and three-body absorption for single-charge exchange. We then include these new elements into an eikonal approximation in order to evaluate differential cross sections for exclusive pion-nuclear single-charge exchange.

2 Propagation of pions through nuclear matter and nuclei

A pion wave traveling through a medium of constant density would contain a phase $\exp(-iV_{opt}t)$ which tells us that pions are removed from the elastic flux at the rate

$$\Gamma = \frac{1}{N}\frac{dN}{dt} = -2ImV_{opt} \ , \qquad (1)$$

which can be rewritten in terms of the pion selfenergy Π, ($2\omega V_{opt} \equiv \Pi$, ω pion energy)

$$Pdt = -\frac{1}{\omega} Im\Pi dt = -\frac{1}{q} Im\Pi dl \qquad (2)$$

where P is the probability per unit time for a reaction to take place and $dl = \frac{q}{\omega} dt$ is the distance travelled by the pion in dt. We can thus interpret $-\frac{1}{q} Im\Pi$ as the probability per unit length for a reaction to take place.

Our idea is to evaluate the pion-nuclear optical potential in infinite nuclear matter as a function of the density and then use it in finite nuclei, substituting at each point the nuclear density by the local density of the nucleus. This is the essence of the local density approximation which we assume in our scheme. However, we include the effects of finite range of the interaction by means of a convolution of the potential with the range of the interaction [3,4].

However, eq. (2) can provide more information than just the total probability of reaction since, as we shall see, we are able to separate the different contributions to $Im\Pi$ and relate them to the probability for different reactions to take place. This is the most important ingredient in our scheme and which allows us to finally evaluate cross sections for the different reactions.

However, the local probability for a certain reaction to take place is only one of the ingredients needed because the pion has to be followed through the nucleus calculating at every moment the probability that a certain reaction takes place. This means that one gets naturally a set of coupled differential equations which require a simultaneous solution. This second part has been done by means of a Monte Carlo computer simulation which has analogies with some of the cascade codes [5,6]. A sufficiently small interval of time is chosen such that the probabilities for the reactions are small compared to unity; a random number is generated which decides which reaction takes place, or whether there is no reaction, according to their corresponding probabilities. In the latter case the pion continues until it finally undergoes some reaction or leaves the nucleus. It is clear that by such a procedure we are taking into account the loss of flux of the pions from the elastic channel which we know is related to the imaginary part of the optical potential. The simulation does this job for us but also tells us into which of the channels the pion has gone: quasielastic, single-charge exchange or absorption. Note that eqs. (1-2) do not contain the elastic scattering, since in infinite matter the pions propagate forward and only the reaction channels are responsible for the loss of forward pion flux. A similar relation to eq. (2) in nuclei is the optical theorem, which relates the total cross section (elastic scattering included) to the imaginary part of the T matrix, $\tilde{\Pi}^R(q, q')$ (not the optical potential) in the forward direction

$$\sigma_{tot} = -\frac{1}{q} Im\tilde{\Pi}^R(q, q) \quad . \qquad (3)$$

It is clear that for our simulation procedure the input provided by eq. (2) is the most appropriate. However, one can not neglect the elastic scattering, but it enters

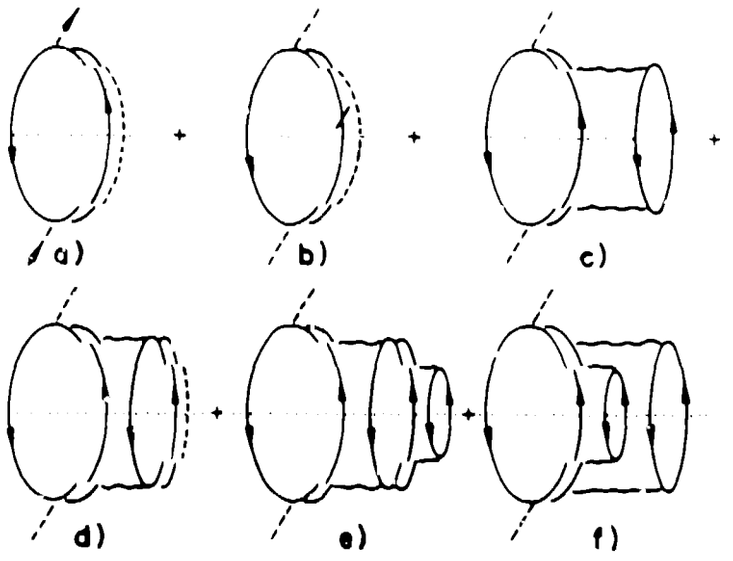


Fig. 1. Model used for the pion-nucleus optical potential, incorporating quasielastic scattering, two-body and three-body absorption. Explanation in the text.

the scheme in an indirect way by distorting the pion waves without any loss of flux. Two approaches have been taken at this point. The first one relies upon the fact that elastic scattering in nuclei is mostly forward and, of course, conserves the pion flux; hence, we take a straight trajectory of the pions between collisions. The second one assumes that between collisions the pion follows a classical trajectory which one evaluates from the real part of a realistic pion-nucleus potential. As we shall see, both procedures lead to somewhat different results at pion energies below $T_\pi = 120$ MeV, but the changes are very small at energies around resonance and above. This gives us confidence that a more realistic quantum mechanical way of dealing with this part would produce similarly small changes.

2.1 Evaluation of the optical potential

Next we show the input for the optical potential. In terms of many-body Feynman diagrams we take the model shown in fig. 1. We evaluate the diagrams in infinite nuclear matter and then construct a local pion-nucleus optical potential by means of the local density approximation, however taking into account the finite range of the interaction, as we mentioned. Now, of course, we want to evaluate and classify all sources of the imaginary part. As we know the sources of imaginary part in the Feynman diagrams appear when the particle lines cut by a straight line are placed on shell. In this way the cuts in diagrams a) and b) would account for quasielastic scattering (including charge exchange). Diagram b) would symbolize the Pauli blocking

correction. In diagram c) the upper part corresponds to two-nucleon pion absorption while the lower cut would account for higher order quasielastic corrections. Similarly the cut shown in diagram d) corresponds again to a higher order quasielastic excitation while the cuts in diagrams e) and f) correspond to three-nucleon pion absorption. The wavy lines stand for the induced interaction which is constructed from π and ρ exchange, modified by the effect of the short range correlations and iterated to all orders of the ph or Δh excitation. Details of the calculation can be seen in ref. 2. The kind of many-body expansion used here bears close resemblance to other successful many-body schemes like the hypernetted chain approach [7] or the planar theory [8].

One comment must be made with respect to the three-body absorption pieces. In diagram f) the part of the interaction which contains pion exchange will be such that the pion will be necessarily off-shell. This is so because a real pion cannot excite a ph in nuclear matter for reasons of energy and momentum conservation. However, diagram e) will have again the pions off-shell in the interaction to the right but the pion in the interaction to the left can be on-shell. This corresponds to a physical process where there is pion quasielastic scattering followed by two nucleon pion absorption. It is clear that such a contribution has to be excluded because it will be automatically generated by the simulation procedure. Hence, what we call three-body absorption contains the contribution from the other pieces of the interaction and the one from off-shell pions.

Another aspect to stress here is the fact that our estimates of four-body absorption [2] and more recent calculations [9] show that these pieces are already smaller than three-body absorption around resonance and somewhat above, and it appears that the many-body approach provides a good convergence at the level of three-particle three-hole excitation at these energies.

2.2 Delta selfenergies

The same model of fig. 1 for the pion selfenergy can be used to construct a model for the Δ selfenergy. Indeed, just erase the external pion lines and the nucleon line of the left in all those diagrams and then we have the corresponding set of Δ selfenergy diagrams. The same cuts discussed above would now be present and would give rise to the imaginary part of the Δ selfenergy. We shall classify these cuts as quasielastic or absorption cuts according to their nature in the corresponding pion selfenergy diagram. As we have discussed before it is clear that some of the cuts correspond to higher order quasielastic scattering and not absorption. The numerical evaluation has been carried out in ref. 10 and we show the results in fig. 2 for the imaginary part of the selfenergy.

We can see that at low energies two-body absorption is dominant, but as the energy increases both the higher order quasielastic part and the three-body absorption get more important and become comparable to the two-body absorption part. Also in fig. 2, the total result is contrasted with the empirical points of the spreading

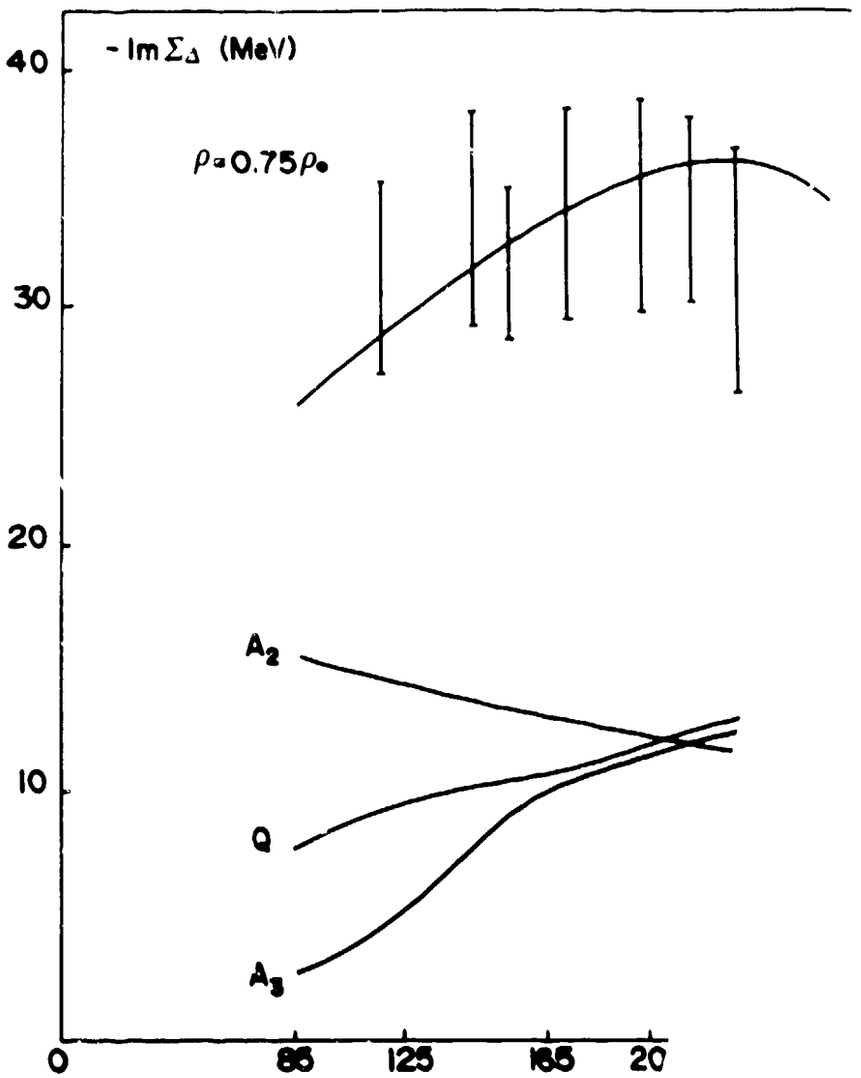


Fig. 2. Imaginary part of the delta selfenergy from ref. 10 split into the contribution from two-body absorption, three-body absorption and higher order quasielastic scattering. The line segments are the empirical determination from ref. 11.

potential of ref. 11 for an effective density of $\rho = 0.75\rho_0$ [12] since the results here are not a function proportional to ρ as assumed in ref. 11. The agreement with the empirical points is rather good, as can be observed in the figure.

3 Simulation procedure

Eq. (2) gave us the probability in the small time Δt for a certain reaction to take place. The procedure outlined above allows for a separation of all different sources of imaginary part of Π and hence provides the separation into the probabilities for each reaction to take place. We start the simulation by generating a random impact parameter at a distance sufficiently far from the nucleus. Then we allow it to proceed forward in steps Δt . At some point it undergoes a reaction or abandons the nucleus. If the pion has been absorbed then it counts as an event for pion absorption. If the reaction has been quasielastic we still have to determine the charge of the outgoing pion, the new direction and the corresponding energy of this pion.

3.1 Determination of the charge

An easy isospin analysis of the structure of the different $\pi N \rightarrow \pi' N'$ amplitudes in the resonance region allows us to write the transition matrix. By assuming a simple scaling of the neutron and proton densities

$$\rho_p/Z = \rho_n/N \tag{4}$$

we can write the transition matrix in terms of the variable x = (N - Z)/A as [13]

$$Q_{\lambda'\lambda} \propto \begin{pmatrix} 5-4x & 1-x & 0 \\ 1+x & 4 & 1-x \\ 0 & 1+x & 5+4x \end{pmatrix} \tag{5}$$

where λ, λ' stand for the π^+, π^0, π^- indices and $Q_{\lambda'\lambda}$ gives the probability that the pion of charge λ becomes a pion of charge λ'. Once again, by generating random numbers and weighting them with these probabilities, we determine the charge of the outgoing pion. The zeros in $Q_{\lambda'\lambda}$ reflect the fact that no double-charge exchange occurs in a single collision in this scheme.

3.2 Determination of the scattering angle

In order to determine the scattering angle we go back and rewrite the $\pi N \rightarrow \pi N$ cross section in the πN CM system as

$$\frac{d\sigma}{d\Omega} \propto Im\left(\frac{1}{\sqrt{s} - M_\Delta + i\Gamma/2 - \Sigma_\Delta(r)}\right)\frac{3\mu^2+1}{2} \tag{6}$$

where μ is cos θ, with θ the angle between the incoming and outgoing pion momenta, and Σ_Δ is the local delta selfenergy.

We proceed in several steps:

1. A random incoming nucleon momentum is chosen from the local Fermi sea in the Lab system.

2. A Lorentz boost is made to transform to the πN CM system.

3. A random angle between the pion momenta is chosen, however, weighted by the angular weight of eq. 6.

4. We make the reverse Lorentz boost to go back to the Lab system.

5. We exclude the quasielastic events in which the outgoing nucleon has a momentum below the local Fermi momentum.

This procedure determines the new direction and energy of the pion in the Lab system. The procedure continues until eventually the pion is either absorbed or it leaves the nucleus with a certain charge, a certain direction and a certain energy, which allows us to determine σ_{abs} or $d^2\sigma/(d\Omega dE)$ for the quasielastic, single-charge exchange and double-charge exchange. By integration over the other variable we can equally determine $d\sigma/d\Omega$ or $d\sigma/dE$.

3.3 S-wave contribution

The resonant mechanisms considered before take into account only the p-wave πN scattering. Below $T_\pi = 100$ MeV the s-wave becomes progressively more important and we include it. The quasielastic contribution is simply taken as $Q^{(s)}dt = \sigma^{(s)}\rho(r)dtq/\omega$, with $\sigma^{(s)}$ the s-wave πN scattering cross section. We also do an isospin analysis based on the experimental phase shifts and construct an energy dependent transition matrix analogous to the one in eq. (5). On the other hand for the true absorption probability we take, as is usual in pionic atoms analysis,

$$A^{(s)}dt = \frac{4\pi}{\omega}\left(1+\frac{\omega}{2M}\right) ImB_0^{(abs)}\rho^2(r)dt \tag{7}$$

with $A^{(s)}$ the probability per unit time for s-wave absorption and M the nucleon mass. $ImB_0^{(abs)}$ is calculated theoretically through methods similar to those exposed before, and we find[14]

$$ImB_0^{(abs)} \approx 0.035 m_\pi^{-4} \ , \tag{8}$$

quite independent of the energy in the range $0 < T_\pi < 100$ MeV.

4 Results for the inclusive pionic reactions

We pass now to a comparison of the results with the experimental data of refs. 15–18.

In fig. 3 we show the results for the total reaction cross section for several nuclei. The agreement with the experimental data is rather good except at low energies

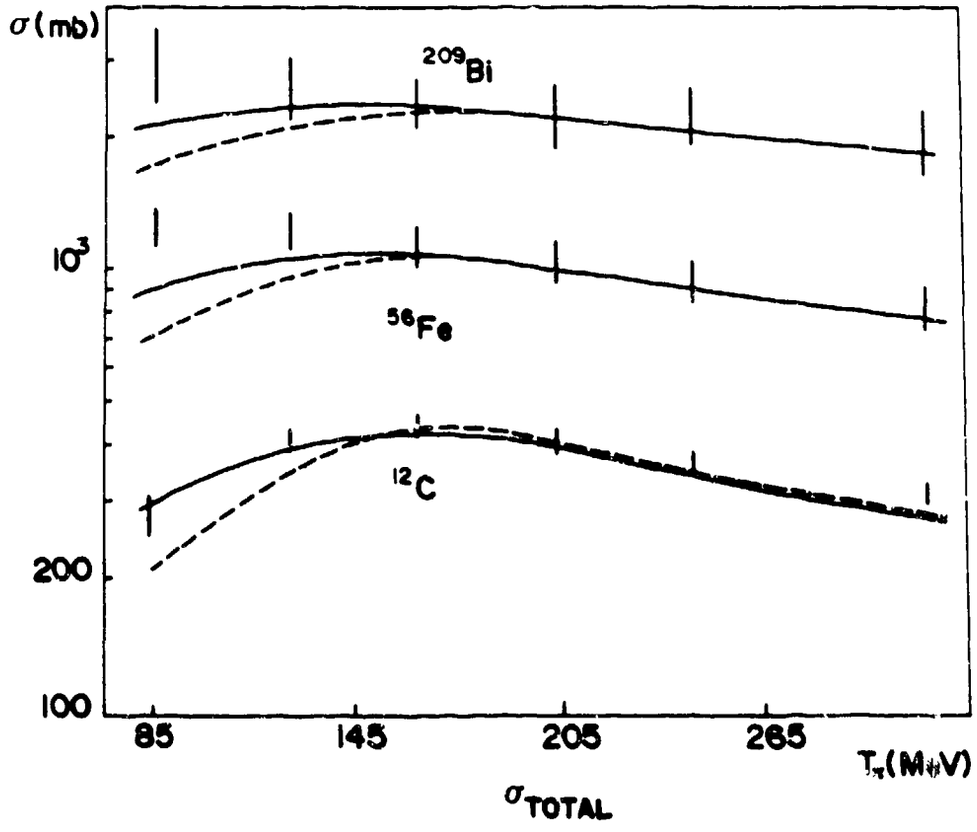


Fig. 3. Total reaction cross section for different nuclei as a function of the energy. The dashed lines assume straight propagation of the pions between collisions. The continuous lines assume classical trajectories between scatterings. Experimental points from ref. 15.

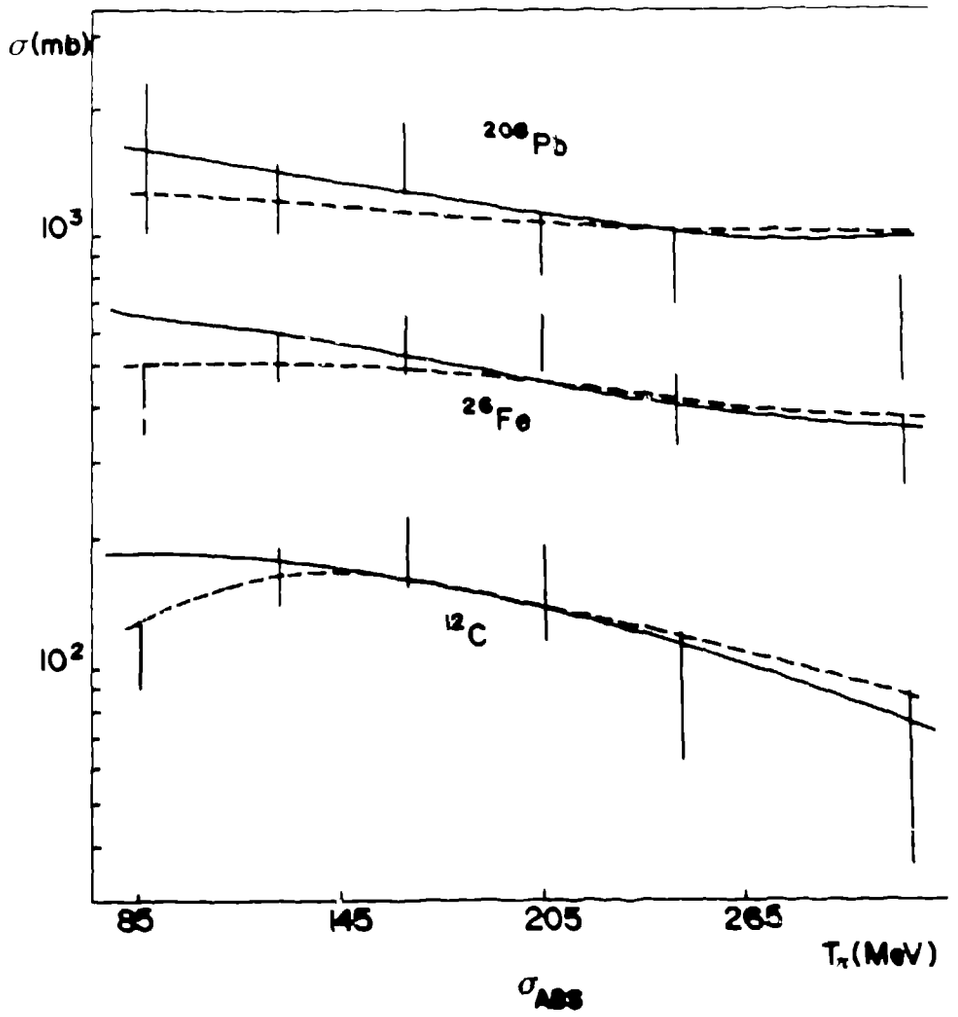


Fig. 4. Absorption cross section. Experimental points from refs. 15.

where we know our approximations of the elastic scattering should be improved. Indeed, we can observe that there are some differences between our results calculated with the straight propagation between collisions or with classical trajectories. These latter results seem to improve the calculation somewhat. For the real part of the optical potential we take a first-order p-wave optical potential arising from ph and Δh excitations [19] plus a s-wave part from ref. 12.

In fig. 4 we show the results for pion absorption with similar features as before and overall good agreement except for the region of low energies.

In fig. 5 we show the results for the integrated quasielastic cross section for two nuclei as a function of the energy. Once again we see that the agreement is rather good. There are large discrepancies with the data of ref. 15 at low energies for ^{56}Fe but the agreement is much better with the new experimental results of ref. 17.

In fig. 6 we show the results for the single-charge exchange integrated cross section

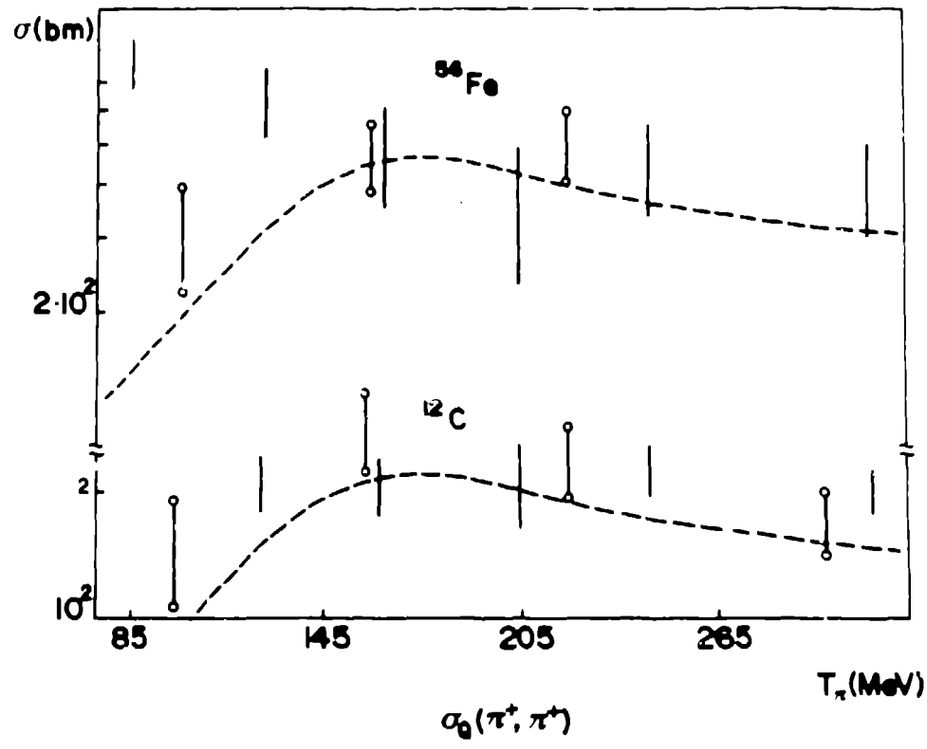


Fig. 5. Quasielastic cross section. Experimental points from refs. 15–17.

for two nuclei. Once again we can observe that the agreement is rather good.

In Fig. 7 we show the results for the double-charge exchange integrated cross section for ^{40}Ca and compare them with the experimental data of ref. 18. The overall agreement is also quite good although at large energies there is more disagreement with the data. However, this is not truly so. Indeed the experimental results are obtained by detecting a π^- in the final state. This experiment thus includes events of pion-induced pion production, essentially $(\pi^+,\pi^-\pi^+)$ events. These cross sections are not negligible compared to those of double-charge exchange without particle production. Indeed, there have been recent measurements at TRIUMF[20] for the $(\pi^+,\pi^+\pi^-)$ reaction on ^{16}O which agree with an earlier calculation of ours[21]. These calculations would give $\sigma \approx 3$ mb for this cross section at $T_\pi = 260$ MeV for ^{40}Ca, bringing the experimental results in closer agreement with the theory.

Finally, in fig. 8 we have selected some results for which there is some discrepancy between our results and the experimental data. These are the angular distributions in ^{208}Pb at $T_\pi = 160$ MeV. Although the theoretical results are in fair agreement with experiment for the quasielastic and (π^+,π^0) reactions, those for the (π^-,π^0) reaction show discrepancies of around a factor two or less. At this point we would like to recall that we have taken the same neutron and proton radius for the nuclear densities in our calculations, an approximation which should not be too good for a heavy nucleus like ^{208}Pb. On the other hand we have taken the Pauli blocking

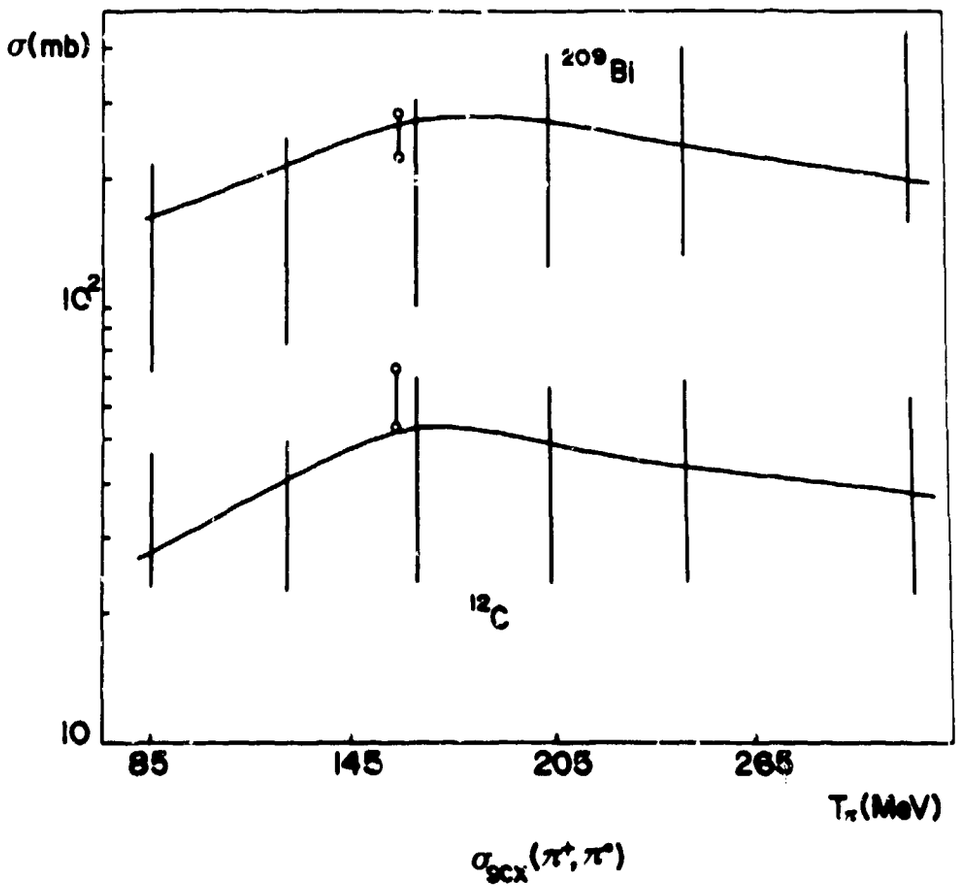


Fig. 6. Single-charge exchange cross section. Experimental results from refs. 16-17.

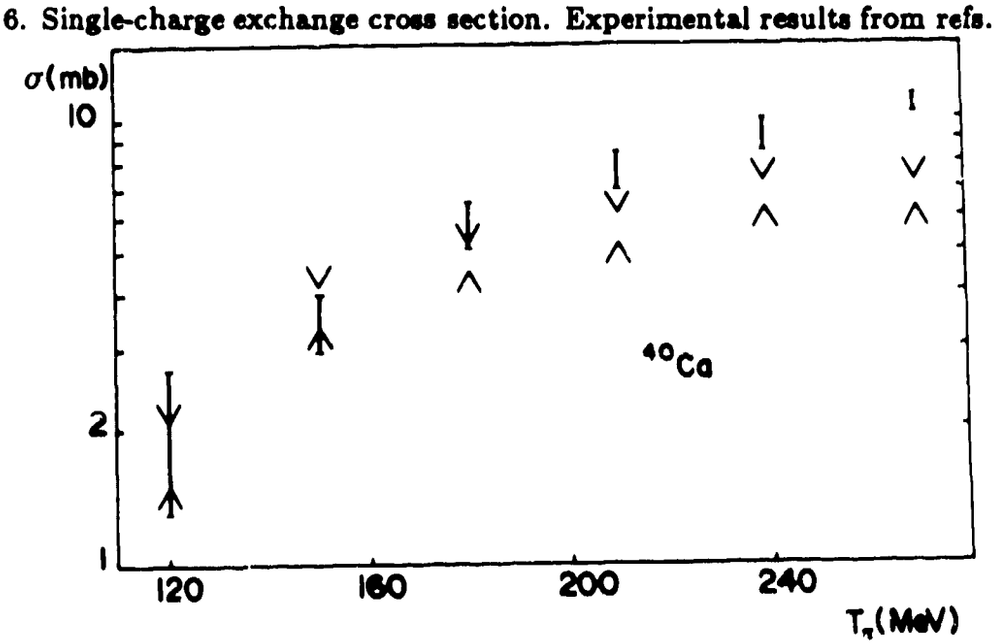


Fig. 7. Double charge exchange cross section. Experimental points from ref. 18. Space between angular symbols: theoretical results with the statistical errors.

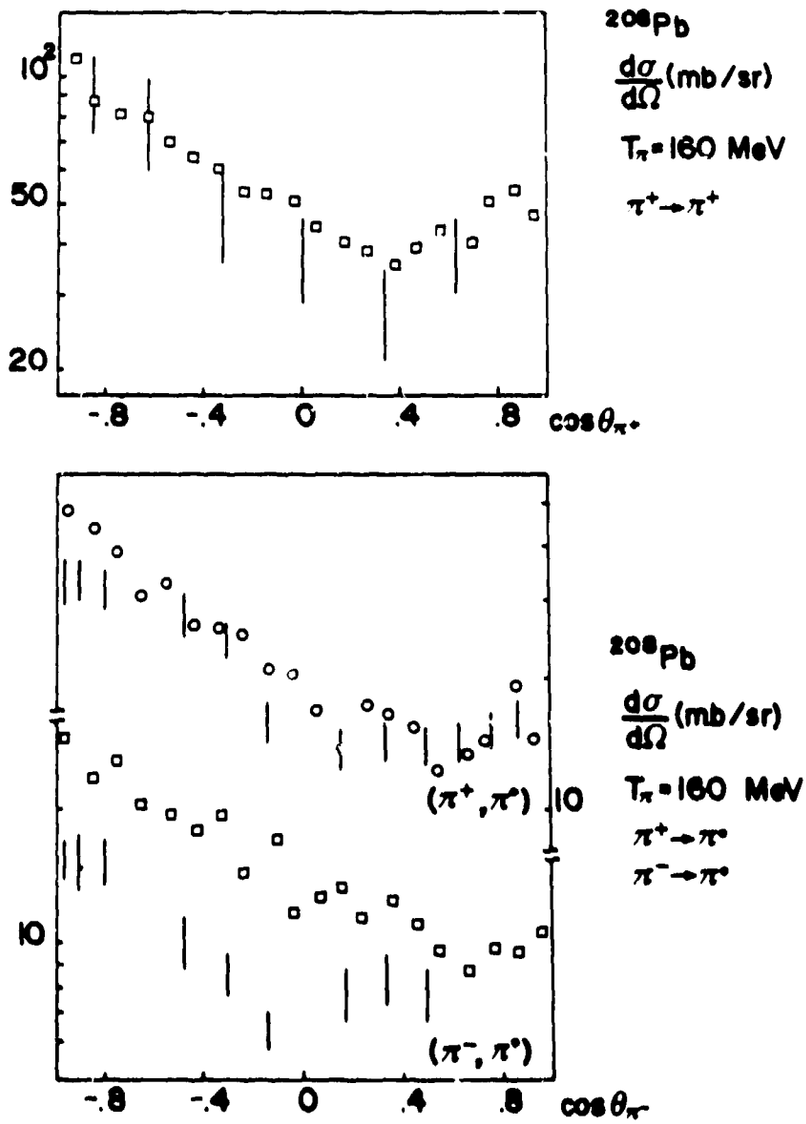


Fig. 8. Angular distributions in quasielastic and single-charge exchange reactions in ^{208}Pb at T_π = 160 MeV. Experimental points from ref. 17.

effect with an average Fermi sea for neutrons and protons. Such details should be incorporated into the scheme for a closer comparison with quasielastic and charge exchange experiments in heavy nuclei where N and Z are quite different.

The overall agreement for different reactions, nuclei and a wide range of energies should be considered as a remarkable achievement of a theory where there are no free parameters.

5 Higher energy exploration

Recent [22] and proposed experiments at energies higher than resonance at LAMPF will bring new information about pion propagation in nuclei. Although in the region of energies around resonance, the key ingredient was the Δ selfenergy, at larger energies where the Δ does not dominate the πN amplitude, a new approach seems inevitable and little has been done theoretically. On the other hand one can envisage that an eikonal approach to the nuclear reactions should be more accurate than at lower energies. At this point one has to raise a warning sign. Glauber theory seems to work remarkably well for some reactions around resonance and even below[23]. This might in principle seem contradictory since one of the essential ingredients in pion-nuclear reactions is pion absorption which affects all nuclear cross sections. However, Glauber theory, which relies upon the elementary πN amplitude alone, does not account for pion absorption which requires at least two nucleons. One of the reasons for this apparent success is the fact that the absorption part of the Δ selfenergy, which increases the Δ width because it incorporates the new $\Delta N \rightarrow NN$ decay channel, partially cancels with the Pauli blocking correction which decreases the Δ width due to some blocking in the $\Delta \rightarrow \pi N$ decay. This appears to be the case in view of the equivalent success of a modified eikonal picture which includes explicitly these effects[24]. However, if one goes to higher energies the Pauli blocking effect will become smaller and eventually negligible while pion absorption mechanisms will still be present. This leads one to a new eikonal picture accounting for pion absorption.

As an example let us look at the single-charge cross section (SCX) in a nucleus. The SCX amplitude in the Glauber approach is written as[25]

$$F(q) = \frac{ik}{2\pi}\int d^2b e^{i\vec{q}\vec{b}} < \psi_{fin}|\sqrt{2}\sum_i \Gamma^{(v)}(\vec{b}-\vec{s}_i)\tau_i^{\pm} \prod_{i\neq i}\left(1-\Gamma^{(s)}(\vec{b}-\vec{s}_i)\right)|\psi_{in}\rangle \qquad (9)$$

where Γ is the profile function [25] for an elementary πN amplitude written as

$$f(q) = f^{(s)}(q) + f^{(v)}(q)\vec{\Theta}\cdot\vec{\tau}\ , \qquad (10)$$

with $\vec{\Theta}$ the isospin matrix for the pions. In the product $\Pi(1-\Gamma)$ we recognize the distortion factor due to quasielastic scattering. Indeed, through some transformations,

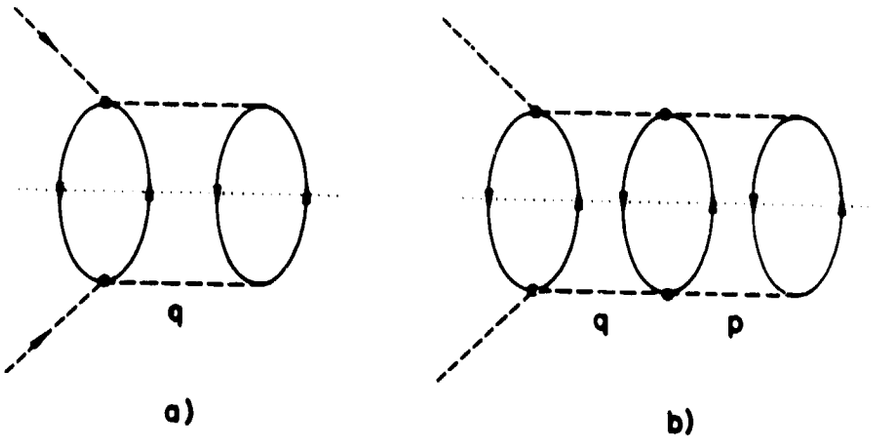


Fig. 9. a) Mechanism for pion absorption through two-particle two-hole excitation mediated by pion exchange; b) same for three-body absorption.

an equivalent way of writing this factor for large values of A is

$$\prod_i (1-\Gamma_i) \approx \exp\left[-\frac{2\pi}{iq_{cm}}\int_{-\infty}^{\infty} dz f(q,q)\rho(\vec{b},z)\right]$$

$$= \exp\left[-i\int_{-\infty}^{\infty} dz \frac{1}{2q_{lab}}\Pi^{(1)}\left(\vec{b},z,q\right)\right] \qquad (11)$$

where $\Pi^{(1)}(\vec{r},q)$ is the lowest order pion selfenergy ($\Pi = -\frac{4\pi\sqrt{s}}{m} f\rho$). This suggests immediately how to include the effect of pion absorption. One simply has to multiply the factor $\Pi(1-\Gamma)$ by

$$F_{abs}(\vec{b}) = \exp\left[-i\int_{-\infty}^{\infty} dz \frac{1}{2q_{lab}}\Pi_{abs}(\vec{b},z,q)\right] \quad . \qquad (12)$$

Since the essential thing done by the pion absorption piece is to remove the absorbed pions from the elastic flux, the crucial thing is to include $Im\Pi^{(2)}_{abs}$ in eq. (12), where $\frac{1}{q} Im\Pi_{abs}$ provides the absorption probability per unit length.

Since we know that two- and three-body absorption is important around resonance[2] one should try to incorporate at least these pieces, which we would find by calculating the imaginary part of the diagrams of fig. 9. A straightforward derivation gives from fig. 9a

$$Im\Pi^{(2)}_{abs} = -\int \frac{d^4q}{(2\pi)^4}|T|^2 ImU(k-q)2D_0(q)^2\frac{f^2}{\mu^2}\vec{q}^{\,2} ImU(q)\theta(k^0-q^0)\theta(q^0) \qquad (13)$$

where T is the πN amplitude, $D_0(q)$ the pion propagator and U(q) the Lindhard function for a ph excitation.

On the other hand one gets for the three body absorption of fig 9b

$$Im\Pi^{(3)}_{abs} = -\int \frac{d^4q}{(2\pi)^4}|T|^2 ImU(k-q)2|\frac{1}{(q^0)^2-\vec{q}^{\,2}-\mu^2-\Pi^{(2)}(q)}|^2 Im\Pi^{(2)}_{abs}(q) \quad (14)$$

where $\Pi^{(2)}(q)$ is the pion selfenergy containing the quasielastic and two-body absorption contribution. Note that in this scheme we are assuming only pions are responsible for the interaction. One could in principle add other ingredients allowed by the $\pi N \rightarrow xN$ virtual transition, but we will concentrate our attention on the pions and on how to separate what we shall call genuine three-body absorption contribution from quasielastic followed by two-body absorption. We will omit the details here, but one can prove that this latter contribution can be obtained from eq. (14) by using the same integral but substituting

$$|\frac{1}{(q^0)^2-q^2-\mu^2-\Pi^{(2)}(q)}| \rightarrow \frac{\pi}{(-)Im\Pi^{(2)}(q)}\delta\left((q^0)^2-\vec{q}^{\,2}-\mu^2\right) \quad . \quad (15)$$

It is clear that since in the Glauber formula, those events where there is a quasielastic scattering will cause the pion to be removed from the elastic flux, one does not have to remove again the pions which undergo quasielastic scattering plus two-body absorption in order to avoid double counting. This means that in $Im\Pi^{(3)}_{abs}(q)$ in the Glauber formula one has to include only the genuine three-body absorption which comes from eq. (14) after subtracting the same integral with the substitution of eq. (15).

The previous discussion solves a formal problem of how to proceed in the many-body method while being consistent with previous assumptions used in the models. On the other hand it is clear that a realistic picture would contain more than just pions being responsible for the exchange of information between particles. In ref. 2, a complete model could be constructed since for resonance dominated processes, objects in the T=1 channel and p-wave couplings (π and ρ) were the essential elements in the exchange. Here, where many partial waves contribute to the πN scattering amplitude and the Δ resonance is no longer dominant, a consistent model is not a trivial task. We have not attempted it, although we believe it to be a problem of prime interest if one wishes to study pionic reactions in that range of energies.

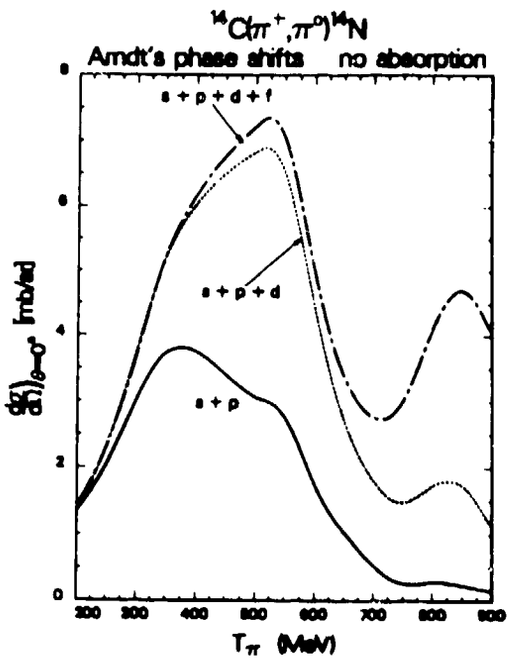
However, it is also interesting to show a first calculation of the effect of the easiest and cleanest piece, two-body absorption mediated by one pion exchange as shown in fig. 9a. We use for such a purpose eq. 13. For the value of $|T|^2$ we have taken the angular average of the same magnitude and thus

$$|T|^2 \rightarrow \frac{1}{4}\sum\overline{\sum}|T|^2 \rightarrow 4\pi\frac{s}{m^2}\sigma \quad (16)$$

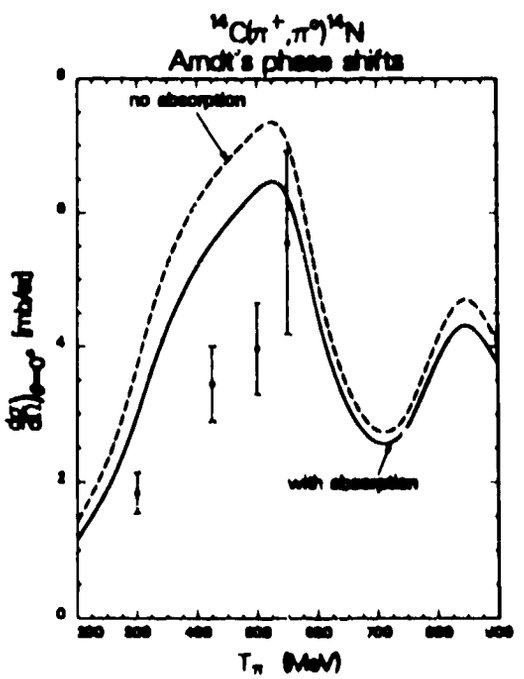
where the average in $|T|^2$ is over initial spin and isospin of the nucleons and the sum is over final spin isospin of the nucleon and isospin of the pion. The quantity σ is then the spin isospin averaged elastic cross section, which is evaluated by using Arndt's phase shifts.[26]

The results that we obtain by using eq. (1) together with the absorption factor, eq. (12) can be seen in figs. 10 and 11.

In fig. 10 we show the importance of including the d wave in the calculation of



SCX cross sections above 300 MeV and of the f wave above 700 MeV for the reaction $^{14}C(\pi^+,\pi^0)$ ^{14}N. Similar effects have been calculated in $^7Li(\pi^+,\pi^0)^7Be$. In fig. 11 we



show the effect of including two-body pion absorption. As discussed before we have not attempted to make a full model for the absorption and we are aware of the many absorption mechanisms beyond the simple two-body absorption mechanism mediated by pion exchange which we have considered. The calculation is however illustrative. We observe that inclusion of this absorption mechanism is responsible for a 20% decrease in the cross section. This is an indication that including other two-body mechanisms, together with the three-body mechanisms, known to be as important as those with two-body at resonance[2], would produce a sizeable decrease of the cross sections which could be estimated at the level around 50%. This would bring the theoretical results closer to the experimental values from refs. 22 and 27, which we also show in the figure. Our results also predict a sharp fall of the differential cross sections beyond $T_\pi = 600$ MeV and a subsequent peak around $T_\pi = 850$ MeV.

6 Conclusions

In the first part of this paper we have shown that a microscopic many-body theory can give proper account of the different reaction cross sections around resonance. The essential ingredient was the Δ selfenergy, particularly its imaginary part and the separation of the imaginary parts coming from different analytical cuts, which were related to the probabilities of reactions in different channels.

As we go above resonance and the Δ does not dominate the reaction one should look for other schemes, different than those used around resonance, in order to study the pion-nuclear reactions. The fact that we are now at higher energies should make an eikonal approach quite reliable. However, unlike proton collisions, in pion-nuclear scattering we have the new channel of pion absorption which is not accounted for in a standard Glauber picture. We have given here a prescription to include it in calculations and have given the first step towards calculating these absorption pieces at higher energies, showing that the SCX experimental data suggest the need for an appreciable effect from pion absorption.

A thorough study of the absorption mechanism, both from the theoretical and experimental point of views looks most interesting at the present time in order to get a good understanding of the different pionic reactions in this new energy region.

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